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Almost No-Scale Supergravity

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Abstract

We construct an explicit 5-dimensional supergravity model that realizes the 'no scale' mechanism for supersymmetry breaking with no unstable moduli. Supersymmetry is broken by a constant superpotential localized on a brane, and the radion is stabilized by Casimir energy from supergravity and massive hypermultiplets. If the standard model gauge and matter fields are localized on a brane, then visible sector supersymmetry breaking is dominated by gravity loops and flavor-violating hypermultiplet loops, and gaugino masses are smaller than scalar masses. We present a realistic model in which the the standard model gauge fields are partly localized. In this model visible sector supersymmetry breaking is naturally gaugino mediated, while masses of the gravitino and gravitational moduli are larger than the weak scale.

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Supersymmetry (SUSY) is arguably the most compelling solution to the hierarchy problem, and SUSY breaking in extra dimensions arguably gives the most attractive solution to the naturalness problems of SUSY [1, 2]. In this Letter, we show that a very simple 5D model can realize SUSY breaking of 'no scale' type, with all moduli fields stabilized. In this model the mass of the gravitino and bulk gravitational moduli are much larger than the scale of supersymmetry breaking in the visible sector. (5D warped compactifications with this feature were recently discussed in Ref. [3].) This elegantly solves the SUSY cosmological problems associated with the gravitino and gravitational moduli. The spectrum is that of gauge-mediated [2] (or radion-mediated [4]) SUSY breaking.

The present model consists of minimal 5-dimensional supergravity (SUGRA) compactified on a S^1/Z_2 orbifold with radius r. We assume that the brane tensions on the orbifold fixed points are small in units of the 5D Planck scale M_5, so that the bulk metric is approximately flat. We assume that there are constant superpotentials localized on one or both orbifold fixed points. This gives rise to a KK spectrum of spin-$\frac{3}{2}$ fermions [5] [1]

$$m_{3/2}^{(n)} = \frac{\pm n + (\delta_0 + \delta_1)/2\pi}{r}, \qquad n = 0, 1, \ldots \tag{1}$$

Here ± 1 is the intrinsic orbifold parity of the mode, and

$$\delta_{0,1} = 2 \tan^{-1} \frac{c_{0,1}}{2M_5^3}, \tag{2}$$

where $c_{0,1}$ are the constant superpotentials localized on the two branes and M_5 is the 5D Planck scale (normalized as defined below). We will assume for simplicity that $c_1 = 0$, and $c = c_0 \neq 0$.

The graviton KK spectrum is not affected by the constant superpotentials, so for $c/M_5^3 \ll 1$, the spectrum is nearly supersymmetric. (We expect $c/M_5^3 \ll 1$ if c arises from gaugino condensation on the brane.) In this case the 4D effective field theory can be written as a supersymmetric effective field theory with SUSY broken spontaneously by the F term of the radion. The effective lagrangian is

$$\begin{aligned} \mathcal{L}_4 = \int & d^4\theta \, \phi^\dagger \phi \left[-3M_5^3 (T + T^\dagger) + \Delta K(T^\dagger, T) \right] \\ & + \left(\int d^2\theta \, \phi^3 \left[c + \Delta W(T) \right] + \text{h.c.} \right), \end{aligned} \tag{3}$$

where $T = \pi r + \cdots$ is the radion chiral multiplet, $\phi = 1 + \theta^2 F_\phi$ is the conformal compensator of 4D SUGRA. This effective theory follows directly from the formulation of 5D SUGRA in terms of $\mathcal{N} = 1$ superfields given in Ref. [6]. We have added a

[1] There is also a single massless spin-$\frac{1}{2}$ fermion, the superpartner of the radion.

radion-dependent Kähler potential and superpotential, which requires additional 5D fields and interactions to be discussed below.

If we neglect ΔW and ΔK, we find that supersymmetry is broken by

$$\langle F_T \rangle = \frac{c^*}{M_5^3}, \tag{4}$$

and the gravitino mass is

$$m_{3/2} = \frac{c}{M_4^2}, \tag{5}$$

where $M_4^2 = M_5^3 2\pi r$ is the 4D Planck scale. The potential vanishes identically, and $\langle F_\phi \rangle = 0$. This is the 'no scale' limit. This limit will certainly not survive quantum corrections, and the radius must be stabilized to obtain acceptable phenomenology. Comparing Eq. (4) with Eqs. (1) and (2), we see that the 4D effective theory is valid if $\langle F_T \rangle \ll 1$.

We now include ΔW and ΔK as perturbations. To linear order,

$$V = -\frac{|c|^2}{M_5^3} \Delta K_{T^\dagger T} - \frac{1}{M_5^3} \left(c^* \Delta W_T + \text{h.c.} \right), \tag{6}$$

where $\Delta W_T = \partial(\Delta W)/\partial T$, *etc.* Higher order terms are suppressed by additional powers of M_5^3. There is no reason for the Kähler and superpotential contributions to the potential to be of the same size, so we expect one or the other to dominate. If the superpotential term dominates, it is easy to see that there is no stable minimum at linear order in ΔW.[2] If the Kähler contribution dominates, there is a stable minimum provided that $\Delta K_{T^\dagger T}$ has a local maximum. We therefore look for stabilization mechanisms that give a nontrivial Kähler potential for the radion.

Assuming that the potential Eq. (6) has a stable minimum, we can cancel the cosmological constant by adding an additional source of SUSY breaking on one of the branes. This adds a positive constant to the right-hand side of Eq. (6). We avoid contact terms between the visible sector and the SUSY breaking sector by assuming that these sectors are localized on different branes. We then obtain

$$\langle F_\phi \rangle = \frac{c^* \Delta K_{T^\dagger T}}{3 M_5^6}. \tag{7}$$

[2]The stabilization mechanism of Ref. [7] makes use of a radius-dependent dynamical superpotential from bulk gaugino condensation. Consistent with the present analysis, this does not lead to SUSY breaking of the no-scale type.

This will give an anomaly-mediated contribution to visible SUSY breaking [8]. As long as

$$\langle \Delta K_{T^\dagger T} \rangle \ll \frac{M_5^3}{2\pi r}, \tag{8}$$

we have $\langle F_\phi \rangle \ll m_{3/2}$, and SUSY breaking masses in the visible sector can be small compared to $m_{3/2}$. We call such a model a 'almost no-scale' model. Eq. (8) is naturally satisfied if the characteristic mass scale of the stabilization dynamics is below M_5.

A natural candidate for Kähler stabilization is Casimir energy [9]. The gravitational contribution to the Casimir energy in this model is [5]

$$V_{\text{grav}} = (-4)\frac{3\zeta(3)}{8\pi^2 L^4}|F_T|^2 + \mathcal{O}(F_T^4), \tag{9}$$

where $L = 2\pi r$. We see that the gravitational contribution to the Casimir energy is attractive, *i.e.* favors small values of the radius.

The KK spectrum of a massive hypermultiplet in this theory can be worked out using the results of Refs. [11]. At each KK level there are 2 states with orbifold parity $+1$ and 2 states with orbifold parity -1:

$$\left(m_0^{(n)}\right)^2 = \left(\frac{\pm n + F_T/2\pi}{r}\right)^2 + m^2, \quad \left(m_{1/2}^{(n)}\right)^2 = \left(\frac{n}{r}\right)^2 + m^2, \quad n = 0, 1, \ldots \tag{10}$$

where ± 1 is the intrinsic orbifold parity. The Casimir energy is then (see *e.g.* Ref. [10])

$$\begin{aligned} V_{\text{hyper}} = (+2)\frac{3}{8\pi^2 L^4}\Big[\tfrac{1}{3}(mL)^2 \mathrm{Li}_1(e^{-mL}) \\ + (mL)\mathrm{Li}_2(e^{-mL}) + \mathrm{Li}_3(e^{-mL})\Big]|F_T|^2 + \mathcal{O}(F_T^4). \end{aligned} \tag{11}$$

The asymptotic behavior is

$$V_{\text{hyper}} \to \begin{cases} (+2)\dfrac{3\zeta(3)}{8\pi^2 L^4}|F_T|^2 + \mathcal{O}(F_T^4) & \text{for } L \ll 1/m, \\ (+2)\dfrac{m^2 e^{-mL}}{8\pi^2 L^2}|F_T|^2 + \mathcal{O}(F_T^4) & \text{for } L \gg 1/m. \end{cases} \tag{12}$$

The asymptotic behavior is easily understood physically. For $L \ll 1/m$ the mass is negligible, and the Casimir energy goes like $1/L^4$ on dimensional grounds. The exponential suppression for $L \gg 1/m$ is the Yukawa suppression of a massive scalar propagating over distances of order L.

If there are 3 or more hypermultiplets, the potential for $L \ll 1/m$ is dominated by the repulsive hypermultiplet contribution, while for $L \gg 1/m$ it is dominated by the attractive gravitational contribution. There is therefore a minimum at $L \sim 1/m$. The fact that massive modes can lead to Casimir stabilization was pointed out by Ref. [10].

The radion mass is $m^2_{\rm radion} \sim \langle F_T \rangle^2 m^4/(16\pi^2 M_4^2)$ and $K_{T^\dagger T} \sim \langle F_T \rangle^2 m^4/(16\pi^2)$. At the minimum the Casimir energy is negative and of order $F_T^2 m^4/(16\pi^2)$. This contribution to the cosmological constant can be cancelled by an additional source of supersymmetry breaking localized on one of the branes. As long as $m \ll M_5$ this can be treated as a perturbation on the analysis above.

We now discuss supersymmetry breaking in the visible sector. As already mentioned, we localize the visible sector and the SUSY breaking sector on different branes to avoid flavor-violating contact terms between these sectors. We must also consider possible flavor-violating contact terms arising from the bulk hypermultiplets. Tree-level contributions from the bulk hypermultiplet can be forbidden by imposing a Z_2 symmetry on the hypermultiplet. The leading couplings of the hypermultiplet to the visible and hidden sectors that cannot be forbidden by symmetries are

$$\Delta\mathcal{L}_5 \sim \int d^4\theta \left[\delta(x^5) \frac{1}{M_5^3} Q^\dagger Q H^\dagger H + \delta(x^5 - \pi r) \frac{1}{M_5^3} X^\dagger X H^\dagger H \right]. \tag{13}$$

These give rise to flavor-violating contact terms in the 4D theory of order

$$\Delta\mathcal{L}_4 \sim \int d^4\theta \, \frac{1}{16\pi^2} \, \frac{1}{M_4^4 L^2} Q^\dagger Q X^\dagger X. \tag{14}$$

It is interesting to consider the possibility that SUSY breaking is communicated to the visible sector by anomaly mediation. However, for the Casimir stabilization considered above, the SUSY breaking from Eq. (14) is always larger than the anomaly-mediated contribution. To see this, we write $\langle \Delta K_{T^\dagger T} \rangle \sim \epsilon m^4$ (where $\epsilon \sim 1/(16\pi^2)$ for Casimir stabilization) and treat L as independent of m. Using $\langle F_\phi \rangle \sim \langle F_T \rangle \epsilon m^4 / M_5^3$ and $\langle F_X \rangle^2 \sim \langle F_T \rangle^2 \epsilon m^4$ we obtain

$$\frac{\Delta m^2_{\rm AMSB}}{\Delta m^2_{\rm loop}} \sim \frac{\epsilon}{16\pi^2} (mL)^4. \tag{15}$$

For Casimir stabilization the right-hand side is of order 10^4, but anomaly mediation may dominate for other radius stabilization mechanisms that give a flatter radion potential. (Because the hypermultiplet loop contribution violates flavor, we need $\Delta m^2_{\rm AMSB} \gtrsim 10^2 \Delta m^2_{\rm loop}$ to suppress flavor-changing neutral currents.)

We instead consider a different mechanism for transmitting SUSY breaking to the visible sector. We assume that the standard-model gauge multiplet is not completely localized, but rather 'leaks out' somewhat into the bulk. $\langle F_T \rangle \neq 0$ then gives rise to a gaugino mass [4] with a suppression factor due to the localization.

We can give an explicit realization of this scenario by assuming that brane-localized kinetic terms dominate over the bulk kinetic term for the gauge zero modes [12]. We have

$$\frac{1}{g_4^2} = \frac{1}{g_{\rm bdy}^2} + \frac{L}{g_5^2} \tag{16}$$

where g_4 is the 4D effective gauge coupling, g_5 is the 5D gauge coupling, and $1/g_{\rm bdy}^2$ is the coefficient of the brane-localized kinetic term. In order for the brane-localized contribution to dominate, we must have

$$\frac{L}{g_5^2} \ll \frac{1}{g_{\rm bdy}^2} \sim \frac{1}{g_4^2} \sim 1. \tag{17}$$

This requires small radius, while we require large radius for sequestering. To see that there is a solution, we must be more careful about numerical factors.

We estimate the size of unknown counterterms by assuming that bulk interactions are strongly coupled at a scale cutoff Λ [13]. This gives

$$\Lambda^3 \sim \ell_5 M_5^3, \qquad g_5^2 \sim \frac{\ell_5}{\Lambda}, \tag{18}$$

where $\ell_5 = 24\pi^3$ is the 5D loop factor. Assuming that massive flavor-violating bulk modes have masses of order Λ, sequestering requires $e^{-\Lambda L} \lesssim 10^{-3}$, so $\Lambda L \gtrsim 7$ is sufficient. We then have

$$\frac{L}{g_5^2} \gtrsim \frac{\Lambda L}{\ell_5} \sim 10^{-2}. \tag{19}$$

The gaugino mass is given by

$$m_{1/2} \sim \frac{\langle F_T \rangle}{g_5^2}, \tag{20}$$

so that

$$\frac{m_{3/2}}{m_{1/2}} \sim \frac{\ell_5}{\Lambda L} \lesssim 100. \tag{21}$$

We see that the gravitino mass can be ~ 100 times larger than the weak scale in this model. This is plausibly enough so that $m_{3/2} > 60$ TeV, in which case the gravitino decays early enough to avoid problems with nucleosynthesis [14].

To avoid flavor-changing neutral currents, we must have $m_{1/2}^2 \gtrsim 10^3 \Delta m_{\rm scalar}^2$, where $\Delta m_{\rm scalar}^2$ is the flavor-violating scalar mass contribution from operators of the form Eq. (14).[3] This is satisfied as long as $\Lambda L \gtrsim m^2/M_4^2$, which is always satisfied.

[3] Gravity loops give flavor-blind contributions to the soft masses that are also of order $\Delta m_{\rm scalar}^2$.

We conclude that this gives an interesting model in which the spectrum is gaugino-mediated, and yet the gravitino is heavy.

We briefly comment on the radiative stability of this model. 4D gravity loops are cut off in the ultraviolet for momenta $p_4 \sim 1/r$, where the extra dimension becomes important. In the 5D theory, gravity loops that contribute to SUSY breaking must connect the two branes, and the contribution from $p_4 \gg 1/L$ is therefore suppressed by $e^{-p_4 L}$. The extra dimension acts as a 'low' cutoff (below M_4) that makes this scenario radiatively stable.

If the gaugino is more strongly localized, we can get a larger value for the ratio $m_{3/2}/m_{1/2}$. The consistency of such a scenario with local 5D supersymmetry is strongly suggested by the fact that we can write an effective theory where the gauge fields are completely localized on an effective-theory brane. An additional hint comes from soliton solutions in higher-dimensional supergravity, which give rise to solutions where $U(1)$ gauge fields are localized on a brane [15].

In a more non-minimal model, there may be additional gravitational moduli, light scalar fields that interact only by Planck-suppressed operators, and which have an exactly flat potential in the SUSY limit. Such fields will also get a mass of order $m_{3/2}$ from contact terms with the hidden sector. If $m_{\rm moduli} \gtrsim 100$ TeV, this solves the Polonyi problem associated with these moduli.

In conclusion, we have presented a simple 5D model that realizes the no-scale mechanism for supersymmetry breaking, with the radion stabilized by Casimir forces. The superpartner spectrum in this model is gaugino mediated, and the main new feature is that the mass of the gravitino and gravitational moduli are heavier than the weak scale, eliminating cosmological problems. More generally, the lesson from this paper (and from Ref. [3]) is that the gravitino mass is controlled by UV physics that is independent from that which gives rise to visible sector SUSY breaking.

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